

Quantum corrections to screening at strong coupling

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Abstract

We compute a certain class of corrections to (specific) screening lengths in strongly coupled nonabelian plasmas using the AdS/CFT correspondence. In this holographic framework, these corrections arise from various higher curvature interactions modifying the leading Einstein gravity action. The changes in the screening lengths are perturbative in inverse powers of the 't Hooft coupling or of the number of colours, as can be made precise in the context where the dual gauge theory is superconformal. We also compare the results of these holographic calculations to lattice results for the analogous screening lengths in QCD. In particular, we apply these results within the program of making quantitative comparisons between the strongly coupled quark-gluon plasma and holographic descriptions of conformal field theory.

Contents

1	Introduction	2
2	Screening masses	3
2.1	Scalar excitations	4
2.2	Vector excitations	6
2.3	Massive spin-2 excitations	7

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3	Quantum corrections to screening	7
3.1	Scalar excitation	10
3.2	Vector excitation	11
3.3	Massive spin-2 excitations	12
4	Comparison with QCD	13
5	Discussion	16
A	Gauge conditions and other functions for eight derivative correction	17
B	Higher derivative corrections to η/s	19

1 Introduction

The AdS/CFT correspondence [1, 2, 3, 4] has proved to be a useful tool in probing strongly coupled physics. The quark gluon plasma [5] formed in heavy ion collisions in the RHIC and LHC may be strongly coupled and it is a useful exercise to see which features can be reproduced using the correspondence [6, 7]. One of the main reasons for this possible connection is the observation that a wide class of holographic theories [8, 9] have a small ratio of shear viscosity (η) to entropy density (s)¹, which is what appears to be needed to explain the RHIC and LHC data [11].

It is now known that while holographic theories describing isotropic plasmas² using Einstein gravity have $\eta/s = 1/4\pi$, this is no longer true when higher curvature corrections are taken into account [13, 14, 15]. While most of the work in the literature has focused on higher curvature corrections to transport coefficients, in this paper we will turn our attention to corrections to screening lengths originally considered in [16]. The screening lengths are defined through spatial correlators of Polyakov loops. When the spatial separation of the loops is large, the fall off of the connected contribution is exponential in separation, namely $e^{-|\mathbf{x}|/\xi}$, where ξ is the screening length. The mass corresponding to the longest screening length is called the mass gap m_{gap} . On the holographic side, this corresponds to the lightest supergravity mode [17, 18, 19, 20] that is exchanged between two strings stretching between the boundary and the horizon. In Einstein gravity, this happens to be the time-time component of the metric. The mass corresponding to the exchange of the axion field is the Debye mass. Depending on which supergravity mode is exchanged, there is a corresponding screening length. Study of pure gluon theories reveal little dependence on N_c so it may be hoped that the differences between the QCD results and the large N_c results are going to be small [21, 22]. The holographic screening masses actually work out to be larger than the lattice results [16] and it could be hoped that finite $\lambda, 1/N_c$ corrections would lead to a better agreement.

The most general higher curvature corrections at leading order are not known in IIB string theory. Only certain special classes of corrections have been worked out [23, 24, 25]. In the large λ and large N_c limit, it will turn out that the known curvature corrections are sufficient to compute the leading perturbative correction to certain screening lengths, namely the mass gap and the screening mass corresponding to a vector and a spin-2 exchange as explained below. In certain supersymmetric theories where one has $\text{AdS}_5 \times M_5$ where M_5 is smooth compact manifold, the leading higher curvature correction is R^4 and gives $1/\lambda^{3/2}$ and $\lambda^{1/2}/N_c^2$ corrections [15, 26]. In certain instances where the

¹It has become a whole new industry to see which fluid in nature has the lowest η/s [10].

²It has been recently claimed that anisotropy can lead to non-universal η/s [12].

dual includes fundamental fields, the leading corrections are $1/N_c$ coming from R^2 terms [27, 28]. As we will argue below, knowledge of the higher curvature terms is sufficient to compute the leading corrections to the screening masses corresponding to the exchange of a graviton. Corrections to screening masses corresponding to the dilaton, axion (Debye mass), C_2 , B_2 are all beyond the scope of current technology (see the discussion in [16]).

In a broader context, the higher derivative modifications of the leading supergravity action will expand the universality class of the dual CFT, by modifying the parameters in the n-point functions of the CFT. In particular, R^3 is a natural term that would appear but only in a non-supersymmetric context [29]. This has been studied in recent toy models [30] where the coupling of the R^3 terms were not restricted to be small. In the context of the present work, we will work in a perturbative framework. A further motivation in considering such corrections is to extend the phenomenology program initiated in [28] and studied further in [31, 32].

An outline of the paper is as follows: In section 2, we review the calculation of the screening masses for scalar, vector and spin-two symmetry channels in the field theory. In section 3, we compute the leading corrections to specific screening masses. In section 4, we compare our results with QCD. Two appendices contain certain details about the calculations used in the paper.

2 Screening masses

We will be interested in the screening masses arising due to exchange of the graviton. The graviton fluctuations can be split into scalar, vector and spin-2 symmetry channels [33]. For convenience, the screening masses in the $\mathcal{N} = 4$ plasma arising from 2-derivative gravity is reproduced [16, 20] in table 1. Here \mathcal{P} , \mathcal{C} and \mathcal{T} represent parity, charge conjugation and time reversal and in the Hilbert space interpretation respectively. \mathcal{CT} corresponds to the Euclidean time reversal.

SUGRA modes	J^{PCT}	m_0	SYM operator
g_{00}	0^{+++}	2.3361	T_{00}
g_{ij}	2^{+++}	3.4041	T_{ij}
a	0^{-+-}	3.4041	$tr(E.B)$
ϕ	0^{+++}	3.4041	\mathcal{L}
g_{i0}	1^{++-}	4.3217	T_{i0}
B_{ij}	1^{+--}	5.1085	\mathcal{O}_{ij}
C_{ij}	1^{--+}	5.1085	\mathcal{O}_{30}
B_{i0}	1^{--+}	6.6537	\mathcal{O}_{i0}
C_{i0}	1^{+--}	6.6537	\mathcal{O}_{3j}
G_a^a	0^{+++}	7.4116	$tr(F^4)$

Table 1: Spectrum [20, 16]. We are interested in g_{00} , g_{0i} and g_{ij} in this paper.

Our goal is to calculate various screening masses in the deconfined phase. Algebraically it works out to be equivalent to working with various components of the graviton in the soliton background as explained in [16]. This will also facilitate a direct comparison of numerics with [19] and [20]. The reason for this equivalence is following: we consider modes of the form $e^{iEt - ik^i z_i}$, where (E, k^i) and (t, z^i) are three dimensional vectors. We get the black hole solution by a double analytic continuation of the coordinates $\tau \rightarrow it$ and $t \rightarrow i\tau$, where τ is the compactified spatial dimension. So the graviton modes in the soliton background, that are of the form e^{iEt} , will correspond to exponentially decaying

modes in the black hole background. With $k^i = 0$, E will correspond to the screening mass. Now we briefly discuss the properties of AdS soliton and calculate various screening masses following the discussion in [19].

The 5-dimensional AdS soliton metric is written as

$$ds^2 = \frac{r^2}{L^2}(f(r)d\tau^2 + \eta_{\mu\nu}dx^\mu dx^\nu) + \frac{L^2}{r^2 f(r)}dr^2, \quad (1)$$

$$\text{with } f(r) = \left(1 - \frac{R_0^4}{r^4}\right).$$

Here $\eta_{\mu\nu}dx^\mu dx^\nu = -dt^2 + dx^2 + dy^2$ and $\eta_{\mu\nu}$ is the 3-dimensional Minkowski metric. Here we choose the notation in which indices a, b etc. will be used for the 5-dimensional spacetime and μ, ν etc. for 3-dimensional Minkowski spacetime. This geometry is constructed by double analytic continuation of a planar AdS black hole and the conical singularity is removed by making the τ coordinate periodic with periodicity $\beta_{soliton} = \frac{\pi L^2}{R_0}$. By double analytic continuation when we get the black hole solution, the periodicity $\beta_{soliton}$ can be related to the temperature of the black hole $T = 1/\beta_{soliton}$.

Now, to determine the screening masses corresponding to a graviton exchange, we solve for the linearized equation of motion for the graviton on the background (1). We write the perturbed metric as

$$g_{ab} = \bar{g}_{ab} + \epsilon h_{ab}, \quad (2)$$

where \bar{g}_{ab} denotes the background metric (1) and ϵh_{ab} is the perturbation with $\epsilon \ll 1$. The background metric is a solution of the five dimensional Einstein-Hilbert action with a negative cosmological constant

$$I = \frac{1}{2\ell_p^3} \int d^5x \sqrt{-g} \left[R + \frac{12}{L^2} \right], \quad (3)$$

where R is Ricci scalar, ℓ_p is the Plank length and L is AdS radius. We insert (2) in this 5-dimensional effective action and collect the terms second order in perturbations; *i.e.*, $\mathcal{O}(\epsilon^2)$. Using this second order lagrangian we find the equations of motion for the perturbations. Now we make the ansatz that the graviton has a solution of the form $h_{ab} = H_{ab}(r)e^{ik \cdot z}$. Here $H_{ab}(r)$ has only radial dependence and vectors z^μ and k^μ are in 3-dimensional spacetime with $k^2 = -M_0^2$; *i.e.*, k is the 3-dimensional momentum vector. We choose the rest frame: $k_\mu = (-M_0, 0, 0)$. Now we can solve the equations of motion by making proper gauge choices that will also ensure that we are classifying the modes into scalar, vector or spin-2 excitations.

2.1 Scalar excitations

The scalar excitations will be related to the diagonal polarization of the graviton and also gives the mass gap in the dual field theory. We begin with the ansatz that the solution is of the following form:

$$\begin{aligned} H_{\tau\tau}(r) &= -\frac{r^2}{L^2}f(r)H_0(r), \\ H_{\mu\nu}(r) &= \frac{r^2}{L^2} \left(\eta_{\mu\nu}a_0(r) + \frac{k_\mu k_\nu}{M_0^2}(b_0(r) + a_0(r)) \right), \\ H_{rr}(r) &= \frac{L^2}{r^2}f^{-1}(r)c_0(r), \\ H_{r\mu}(r) &= ik_\mu d_0(r). \end{aligned} \quad (4)$$

Here $\eta_{\mu\nu}$ is the metric on the Minkowski section of the (1). Choosing $k^\mu = (M_0, 0)$, we get the following simplified form of the perturbations

$$\begin{aligned} h_{\tau\tau} &= -\frac{r^2}{L^2}f(r)H_0(r)e^{-iM_0t}, & h_{tt} &= \frac{r^2}{L^2}b_0(r)e^{-iM_0t}, \\ h_{ii} &= \frac{r^2}{L^2}a_0(r)e^{-iM_0t}, & h_{rr} &= \frac{L^2}{r^2}f^{-1}(r)c_0(r)e^{-iM_0t}, \\ h_{rt} &= -iM_0d(r)e^{-iM_0t}, \end{aligned} \quad (5)$$

where now i represents the spatial coordinates of the Minkowski spacetime. As explained in [19], to solve the equations of motion consistently, we also require

$$\begin{aligned} a_0(r) &= \frac{H_0(r)}{2}, \\ c_0(r) &= \frac{-2R_0^4}{3r^4 - R_0^4}H_0(r) \quad \text{and} \\ d_0(r) &= \frac{r^2}{2L^3M_0^2}\frac{db_0(r)}{dr} - \frac{r^2R_0^4}{(3r^4 - R_0^4)L^3M_0^2}\frac{dH_0(r)}{dr} \\ &\quad - \frac{12r^4R_0^4}{(3r^4 - R_0^4)^2L^2M_0^2}H_0(r). \end{aligned} \quad (6)$$

These conditions are obtained by adopting the following strategy to solve the equations of motion:

- First, we make a choice of the gauge by defining $a_0 = H_0(r)/2$.
- By looking at the equations of motion of h_{tr} , we see that it is satisfied if we choose $c_0(r) = -2R_0^4H_0(r)/(3r^4 - R_0^4)$.
- Using these $a_0(r)$ and $c_0(r)$ we simplify the other equations of motion. If the equation for h_{xx} is subtracted from equation for $h_{\tau\tau}$, we find the expression (6) for $d_0(r)$.
- Now by simplifying the rest of the equations of motions, we get the following differential equation

$$r(r^4 - R_0^4)\frac{d^2H_0(r)}{dr^2} + (5r^4 - R_0^4)\frac{dH_0(r)}{dr} + \frac{r(64r^2R_0^8 + L^4M_0^2(-3r^4 + R_0^4)^2)}{(-3r^4 + R_0^4)^2}H_0(r) = 0. \quad (7)$$

Here, we point out that to solve the equations of motions we do not need to fix $b_0(r)$. In the beginning we could have fixed $b_0(r)$ and then the rest of the conditions would have followed accordingly.

The differential equation (7) is solved for the value of M_0 such that the solution satisfies certain boundary conditions. The required boundary conditions are the following: first, we impose the condition that $H_0(r)$ is finite at the horizon. This condition can be implemented by simply solving (7) near the horizon. We find that close to $r = R_0$, the solution should behave like

$$H_0(r) = 1 - \frac{(16R_0^2 + L^4M_0^2)}{4R_0^3}(r - R_0). \quad (8)$$

The second boundary condition is that $H_0(r)$ should fall off asymptotically as $1/r^4$. This condition comes from the fact that metric fluctuations should fall off precisely with the rate to yield a non-vanishing stress-energy tensor $\langle T_{ab} \rangle$ in the dual field theory [34]. So asymptotically we expect that

$$H_0(r) = \frac{L^4C_1}{r^4} + \frac{L^6C_2}{r^6} + \frac{L^8C_3}{r^8} \dots, \quad (9)$$

where C_1, C_2, C_3 etc. are dimensionless constants.

To solve equation (7) numerically we express (7) and (8) in dimensionless variables u and \mathcal{M}_0 , where $r = uR_0$ and $\mathcal{M}_0 = L^2 M_0 / R_0$. Now we solve the differential equations iteratively and find the numerical value of \mathcal{M}_0 such that the solution satisfies both the boundary conditions. This is the shooting method. We find that screening mass for the scalar symmetry channel, in units of temperature, is

$$M_0 = 2.336\pi T, \quad (10)$$

which is consistent with [19, 20].

2.2 Vector excitations

To study the vector and spin-2 excitations, we start with the following ansatz for the graviton polarizations

$$H_{ab} = \varepsilon_{ab} \frac{r^2}{L^2} H(r). \quad (11)$$

Here, ε_{ab} is a constant polarization tensor and satisfies

$$\varepsilon_{\tau a} = \varepsilon_{ra} = 0 = \varepsilon_{a\mu} k^\mu \quad \forall a, \quad (12)$$

where k^μ is the momentum vector in the Minkowski spacetime. Further, to separate the vector excitations, we impose the condition that the non-vanishing components of the polarization tensor take the following form

$$\varepsilon_{\mu\tau} = \varepsilon_{\tau\mu} = v_\mu, \quad \text{with } k \cdot v = 0 \text{ and } v \cdot v = 1. \quad (13)$$

As we have chosen $k^\mu = (M_{v0}, 0, 0)$, for convenience, we can choose $v^\mu = (0, 1, 0)$. We write $H(r) = H_{v0}(r)$ and with these choices, the only non-zero perturbation is

$$h_{\tau x} = \frac{r^2}{L^2} H_{v0}(r) e^{-iM_{v0}t}. \quad (14)$$

Just to clarify the notation, we point out that in M_{v0} and $H_{v0}(r)$, the subscript ‘ v ’ is for vector excitation and the subscript ‘ 0 ’ is to indicate that we are working with Einstein gravity in (3). We find that the equation of motion for this perturbation is

$$r \frac{d^2 H_{v0}(r)}{dr^2} + 5 \frac{dH_{v0}(r)}{dr} + \frac{L^4 M_{v0}^2 r}{r^4 - R_0^4} H_{v0}(r) = 0. \quad (15)$$

Now this equation of motion is solved numerically for M_{v0} such that the solution satisfies proper boundary conditions. We find that close to the horizon $r = R_0$, the solution behaves like

$$H_{v0}(r) = (r - R_0). \quad (16)$$

The second boundary condition is that, similar to (9), the asymptotic solution should be

$$H_{v0}(r) = \frac{L^4 C_1}{r^4} + \frac{L^6 C_2}{r^6} + \frac{L^8 C_3}{r^8} \dots \quad (17)$$

We can again express the equations (15) and (16) in dimensionless variables and solve the equation of motion numerically. We find that in units of temperature, the screening mass for vector channel is

$$M_{v0} = 4.322\pi T, \quad (18)$$

in agreement with [19, 20].

2.3 Massive spin-2 excitations

The remaining graviton fluctuations, g_{ij} , describe the massive spin-2 excitations on field theory side. We again start with the ansatz (11) (with $H(r) = H_{s0}(r)$) and the polarization tensor satisfies the conditions (12). We impose the tracelessness condition:

$$\eta^{\mu\nu}\varepsilon_{\mu\nu} = 0. \quad (19)$$

Thus, now (11) will describe two independent modes, one of which will be off-diagonal, *i.e.*, $\varepsilon_{xy} = \varepsilon_{yx} = 1$ and otherwise $\varepsilon_{ab} = 0$. Another mode will be the diagonal and traceless with $\varepsilon_{xx} = -\varepsilon_{yy} = 1$, and $\varepsilon_{ab} = 0$ otherwise. For convenience, we will calculate the screening mass for the off-diagonal polarization which does not mix with other polarizations. In this case, the only non-zero perturbations will be

$$h_{xy} = \frac{r^2}{L^2} H_{s0}(r) e^{-iM_{s0}t}, \quad (20)$$

and the equation of motion for it is

$$r(r^4 - R_0^4) \frac{d^2 H_{s0}(r)}{dr^2} + (5r^4 - R_0^4) \frac{dH_{s0}(r)}{dr} + rL^4 M_{s0}^2 H_{s0}(r) = 0. \quad (21)$$

We find that close to the horizon $r = R_0$, the solution is given by

$$H_{s0}(r) = 1 - \frac{L^4 M_{s0}^2}{4R_0^3} (r - R_0). \quad (22)$$

The other boundary condition is that the solution has asymptotic behavior similar to (9). Now we solve the equation of motion (21) numerically for M_{s0} and get [19, 20]

$$M_{s0} = 3.404\pi T. \quad (23)$$

3 Quantum corrections to screening

As we have explained, we are interested in finding out the screening mass in the Yang-Mills plasma corresponding to a graviton exchange in the Polyakov loop correlator on the gravity side. This entails expanding the IIB low energy effective action to quadratic order in fluctuations. In the 10-dimensional action only the metric and the 5-form flux are turned on. This simplifies the problem enormously since we do not have to worry about terms such as RH^2 where H is a combination of the RR and NSNS 3-form field strengths. One problem that we could face is that there could be mixing between the metric fluctuation and the scalar or the gauge field. Since we do not know these terms at higher derivative order completely, our results could be incomplete. Now we will argue that such mixing will not arise. Firstly, we assume that since we are in a perturbative approximation, there is no level crossing, *i.e.*, the lightest mode remains the lightest mode and so on. When we Fourier expand, each mode comes with a factor of $e^{-iM_i t}$ where M_i is the eigenvalue corresponding to that mode. Thus integration over time is only going to allow for mixing at quadratic order between degenerate modes. Looking at table 1, it is thus clear that the only problematic terms could arise due to the mixing of the metric and the dilaton-axion at quadratic order which would effect the result for the spin-2 exchange. Firstly note that the dilaton is constant to leading order. It gets sourced by the leading higher derivative correction at $O(\gamma)$. Thus any fluctuation should also be at least $O(\gamma)$.

Thus plugging this into the action would lead to contributions at least at $O(\gamma^2)$. Since there is an $SL(2, \mathbb{Z})$ symmetry, we would also conclude that the axion would contribute at $O(\gamma^2)$. We can do somewhat better than this. Since there is an $SL(2, \mathbb{Z})$ symmetry in IIB string theory, a term of the form $\partial^m h \partial^n \phi$ (where h and ϕ denote metric and dilaton fluctuations around the background) would necessitate the existence of $\partial^m h \partial^n a$ where a is an axion-fluctuation. The 10-dimensional origin of the latter term could be of the form $C_{abcd} C^{aec}{}_f C^{bd}{}_{eg} \nabla^f \nabla^g C^{(0)}$. However such terms cannot arise in perturbation theory due to parity conservation. As the problematic terms alluded to above will not arise and we need to simply focus on the curvature corrections. At this point, we must point out that although we are guided by type IIB string theory, the phenomenological program in [28] is more general and we expect corrections starting at four derivative order with unknown mixings between various fields. Nevertheless, if we work in a perturbative approximation, our comments above will also apply to the general case.

Now the most well known curvature correction in IIB is the W^4 term with a specific contraction between 4 Weyl tensors dictated by supersymmetry. In general, we may expect that due to addition of fundamental matter, the corrections would begin at W^2 order and in a non-supersymmetric plasma there will also be W^3 terms and a more general set of W^4 terms. We will not consider terms involving covariant derivatives of the Weyl tensor in the non-supersymmetric case—this is an assumption for simplicity. The various tensor contractions of the Weyl tensors are shown in table 2. In the next section we will show in detail the analysis with W^4 which is relevant for the screening masses in the $\mathcal{N} = 4$ plasma.

	W_i^4 terms	W^3 terms	W^2 terms
1.	$W^{wvrs} W_{wv}{}^{tu} W_{rmt}{}^n W_{snu}{}^m$	$W^{rstu} W_{rt}{}^{vw} W_{svuw}$	$W_{rstu} W^{rstu}$
2.	$W^{wvrs} W_{wvr}{}^t W_s{}^{umn} W_{tmun}$		
3.	$W^{wvrs} W_{wr}{}^{tu} W_{vt}{}^{mn} W_{smun}$		
4.	$W^{wvrs} W_{wtr}{}^u W_v{}^{mtn} W_{snum}$		
5.	$W^{wvrs} W_{wtr}{}^u W_{vms}{}^n W^{tm}{}_{un}$		

Table 2: Independent contractions of Weyl tensors in five dimensions

Since we wish to consider general curvature corrections as in table 2, it is instructive to review the supersymmetric W^4 case first in detail to set out the procedure for the remaining cases. We begin with the 5-dimensional effective action in Einstein frame including the eight derivative correction term³ W_s^4

$$I = \frac{1}{2\ell_p^3} \int d^5x \sqrt{-g} \left[R + \frac{12}{L^2} + \gamma L^6 W_s^4 \right], \quad (24)$$

where dimensionless coupling constant [15] $\gamma = \zeta(3)/8\lambda^{3/2} + \sqrt{\lambda}/384N_c^2 \ll 1$ with λ being the 't Hooft coupling⁴ and W_s^4 is given by

$$W_s^4 = W^{hmnk} W_{pmnq} W_h{}^{rsp} W_{rsk}^q + \frac{1}{2} W^{hkmn} W_{pqmn} W_h{}^{rsp} W_{rsk}^q. \quad (25)$$

This form of the corrections appears in the supergravity action and is equal to $(1/2W_1^4 + W_4^4)$, where W_i^4 is the i^{th} -contraction of four Weyl tensors in Table 2.

³These terms originate from the dimensional reduction of the 10-dimensional action consisting of the metric and the five-form flux [35, 15, 26].

⁴Note that if $\lambda = 6\pi$, $N_c = 3$ corresponding to $\alpha_s = 0.5$, then $\gamma \approx 0.003$.

We write the soliton solution as

$$ds^2 = \frac{r^2}{L^2}(f(r)d\tau^2 + \eta_{\mu\nu}dx^\mu dx^\nu) + \frac{L^2}{r^2 g(r)}dr^2, \quad (26)$$

with

$$\begin{aligned} f(r) &= f_0(r)(1 + \gamma f_1(r)), \\ g(r) &= f_0(r)(1 + \gamma g_1(r)) \\ \text{and } f_0(r) &= \left(1 - \frac{R_0^4}{r^4}\right). \end{aligned} \quad (27)$$

We plug in this solution into the action (24) and find equations of motion for $f_1(r)$ and $g_1(r)$. After solving these coupled differential equations we find that

$$\begin{aligned} g_1(r) &= \frac{C_1 r^{12} + 360 r^4 R_0^{12} - 285 R_0^{16}}{r^{12}(r^4 - R_0^4)}, \\ f_1(r) &= \frac{C_1 r^{12} + 120 r^4 R_0^{12} - 45 R_0^{16}}{r^{12}(r^4 - R_0^4)} + C_2, \end{aligned} \quad (28)$$

where C_1 and C_2 are constants. To fix these constants we impose following conditions:

- We demand that the horizon is still at $r = R_0$. So $g_1(r)$ and $f_1(r)$ should be regular at $r = R_0$, *i.e.*, the numerator of the expression for $g_1(r)$ should vanish at $r = R_0$. We get

$$\begin{aligned} C_1 &= -75 R_0^4 \\ \text{and now } g_1(r) &= \frac{-15 R_0^4 (5r^8 + 5r^4 R_0^4 - 19 R_0^8)}{r^{12}}, \\ f_1(r) &= \frac{-15 R_0^4 (5r^8 + 5r^4 R_0^4 - 3 R_0^8)}{r^{12}} + C_2. \end{aligned} \quad (29)$$

Here we notice that by choosing the given C_1 , the stated condition is satisfied by both $g_1(r)$ and $f_1(r)$.

- To fix C_2 we use the following condition: to find black hole solution of the action (24) we just need to do the double analytic continuation in the metric (26); *i.e.*, $\tau \rightarrow it$ and $t \rightarrow i\tau$. In this solution, we can extract the background metric for the dual gauge theory by going to the asymptotic limit giving

$$ds^2 = -f(\infty)dt^2 + d\tau^2 + dx^2 + dy^2. \quad (30)$$

To have the speed of light to be one in the dual gauge theory, one requires $f(\infty) = 1$ and this fixes the constant $C_2 = 0$.

Having found the soliton solution, we remove the conical singularity by compactifying the τ coordinate to be periodic and this periodicity $\beta_{soliton}$ will be related to the temperature of the dual field theory

$$T = \frac{1}{\beta_{soliton}} = \frac{R_0}{L^2 \pi} (1 + 15\gamma). \quad (31)$$

Now to calculate the free energy of the dual gauge theory, we use standard path integral technique in which we identify the Euclidean action (I_{euc}) of the bulk gravity with the ratio of the free energy

and temperature (w/T) of the dual field theory. To render the action finite, following [36], we use background subtraction to compute the free energy:

$$w = -\frac{\pi^4 L^3 T^4}{2\ell_p^3}(1 + 15\gamma). \quad (32)$$

leading to the entropy density:

$$s = -\frac{\partial w}{\partial T} = \frac{2\pi^4 L^3 T^3}{\ell_p^3}(1 + 15\gamma). \quad (33)$$

3.1 Scalar excitation

Having studied the thermodynamic properties of this perturbed solution we now turn to computing the correction to the mass gap. We begin with

$$g_{ab} = \bar{g}_{ab} + \epsilon h_{ab}, \quad (34)$$

where $\epsilon \ll 1$. We insert (34) in the effective action (24) and find the terms of the lagrangian that are second order in ϵ and use these terms to find equations of motion for perturbations. Now we make the ansatz that the solution is of the form $h_{ab} = H_{ab}e^{-ik \cdot z}$, where k and z are three-dimensional vectors in the Minkowski section of metric (26). Here $z^\mu = (t, x, y)$ and k is momentum vector with $k^2 = -M^2$. As we want to find the correction to the mass gap, we expect that up to first order $M = M_0 + \gamma M_1$, where M_0 is screening mass for the scalar excitation given by (10) and M_1 is the first order correction. We further make the ansatz that solution is of the form:

$$\begin{aligned} H_{\tau\tau}(r) &= -\frac{r^2}{L^2}f(r)H(r), \\ H_{\mu\nu}(r) &= \frac{r^2}{L^2}\left(\eta_{\mu\nu}a(r) + \frac{k_\mu k_\nu}{M^2}(b(r) + a(r))\right), \\ H_{rr}(r) &= \frac{L^2}{r^2}f^{-1}(r)c(r), \\ H_{r\mu}(r) &= ik_\mu d(r). \end{aligned} \quad (35)$$

For convenience, we choose the frame where $k^\mu = (M, 0, 0)$ so that the solution is of the form $h_{ab} = H_{ab}e^{-iMt}$. These expressions are similar to (4) but now all the terms have first order corrections in γ . So we write

$$\begin{aligned} H(r) &= H_0 + \gamma H_1, & a(r) &= a_0(r) + \gamma a_1(r), \\ b(r) &= b_0(r) + \gamma b_1(r), & c(r) &= c_0(r) + \gamma c_1(r), \\ d(r) &= d_0(r) + \gamma d_1(r), & M &= M_0 + \gamma M_1, \end{aligned} \quad (36)$$

where H_0, a_0, b_0, c_0 and d_0 are zeroth order solutions. The perturbations are of the following form

$$\begin{aligned} h_{\tau\tau} &= -\frac{r^2}{L^2}f(r)H(r)e^{-iMt}, & h_{tt} &= \frac{r^2}{L^2}b(r)e^{-iMt}, \\ h_{\mu\mu} &= \frac{r^2}{L^2}a(r)e^{-iMt}, & h_{rr} &= \frac{L^2}{r^2}f^{-1}(r)c(r)e^{-iMt}, \\ h_{rt} &= -iMd(r)e^{-iMt}. \end{aligned} \quad (37)$$

The zeroth order terms of the equations of motion ($\mathcal{O}(\gamma^0)$) give the differential equation (7). The first order terms of equations of motion will be functions of $H_0(r)$, $H_1(r)$, $a_1(r)$, $b_1(r)$, $c_1(r)$, $d_1(r)$, M_0 , M_1 and their derivatives. Now we explain how we solve the first order equations of motion for perturbations:

- First, we eliminate the higher derivatives of $H_0(r)$ in the first order equations of motion ($\mathcal{O}(\gamma)$) by using the leading order equation of motion (7) and its derivatives.
 - Similar to the discussion in section 2, we fix the gauge by choosing $a_1(r) = H_1(r)/2$.
 - Now, we find the expression for $c_1(r)$ by demanding that the equation of motion for h_{tr} is satisfied. We have given the expression for $c_1(r)$ in the appendix A. It can be seen that the functional dependence of $c_1(r)$ on $H_1(r)$ is similar to the dependence of $c_0(r)$ on $H_0(r)$ in (6). There are extra terms in expression of $c_1(r)$ that depend on $H_0(r)$ and these terms come from the correction γW_s^4 .
 - After using this expression for $c_1(r)$, we find that the equation for h_{tt} is also satisfied. Now we use the equation for the perturbation h_{rr} and find the functional form of $d_1(r)$. We give the explicit expressions in appendix A.
 - After using this form of $d_1(r)$, the equation for h_{rr} is satisfied and from rest of the equations of motion we get the differential equation for $H_1(r)$. Similar to the leading order solution, we fix the gauge by fixing either $a_1(r)$ or $b_1(r)$ and rest of the consistency conditions follow accordingly. We also notice that the differential equation for $H_1(r)$ (equation (64) in the appendix A) is similar to (7), but with the leading order solution $H_0(r)$ behaving as source.
 - Equations (7) and the equation for $H_1(r)$ are coupled differential equations. Now, knowing the value of M_0 from (10), we can solve for M_1 . The boundary conditions for $H_0(r)$ is given by (8) and (9). Similar to the boundary condition for $H_0(r)$, we solve the equation of motion for $H_1(r)$ (equation (64)) near the horizon $r = R_0$. $H_1(r)$ near the horizon is given by (65). Similar to $H_0(r)$ in (9), the second boundary condition for the solution is that asymptotically $H_1(r)$ falls off as $1/r^4$.
- Finally, using the shooting method we find

$$M = \pi T(2.336 - 139.514\gamma). \quad (38)$$

A similar analysis can be done for all the perturbations given in Table 2. We have show the results for these terms in the Tables 3, 4, 5 and 6.

3.2 Vector excitation

In this section, we calculate the correction to the screening mass for the vector channel. We start with a perturbed metric of the form (34). We insert this metric into the effective action (24) and find the equations of motion for perturbations by singling out the terms $\mathcal{O}(\epsilon^2)$. Now we make the ansatz that $h_{ab} = H_{ab}e^{-ik \cdot z}$, where k and z are three-dimensional vectors in the Minkowski metric (26) with $k^2 = -M_v^2$. We further expect that $M_v = M_{v0} + \gamma M_{v1}$, where M_{v1} is first order correction to the screening mass. We further make the ansatz that the solution is of the following form

$$H_{ab} = \varepsilon_{ab} \frac{r^2}{L^2} H_v(r). \quad (39)$$

Here ε_{ab} is a constant polarization tensor and satisfies (12) and (13). The only non-zero perturbation is

$$h_{\tau x} = \frac{r^2}{L^2} H_v(r) e^{-iM_v t}. \quad (40)$$

For this perturbation, we assume that $H_v(r) = H_{v0}(r) + \gamma H_{v1}(r)$, where H_{v0} is the solution of differential equation (15) and H_{v1} is the first order correction.

To find the equation of motion for $H_{v1}(r)$, we plug in (40) in the equation of motion for the perturbation $h_{\tau x}$. After inserting $M_v = M_{v0} + \gamma M_{v1}$, we expand this equation of motion in powers of γ . We see that the zeroth order terms in γ will give us the equation of motion for $H_{v0}(r)$, *i.e.*, equation (15). We can use this zeroth order equation of motion to eliminate derivatives of $H_{v0}(r)$ from terms first order in γ . Finally, we get equation of motion for $H_{v1}(r)$ that is given in the appendix A (equation (66)).

The equation of motion for $H_{v1}(r)$ and $H_{v0}(r)$ are coupled differential equations. Now we can find M_{v1} by using the value of M_{v0} (equation (18)) and solving this set of coupled differential equations such that $H_{v1}(r)$ satisfies proper boundary conditions. Using (16) we can find the solution for $H_{v1}(r)$ near the horizon and it is given by (67).

The second boundary condition for $H_{v1}(r)$ is that it falls off asymptotically similar to $H_{v0}(r)$, *i.e.*, a function similar to (17). Now we can solve the coupled differential equations (15) and (66) for M_{v1} by converting it into dimensionless parameters (see discussion before (38)). Finally, we find that the screening mass for the vector channel is

$$M_v = \pi T(4.322 - 398.354\gamma). \quad (41)$$

3.3 Massive spin-2 excitations

We start with $h_{ab} = H_{ab}e^{-ik \cdot z}$. We also assume that $M_s = M_{s0} + \gamma M_{s1}$, where M_{s1} is first order correction to the screening mass. Here M_{s0} is given by (26). Further, we make the ansatz that the solution is of the following form

$$H_{ab} = \varepsilon_{ab} \frac{r^2}{L^2} H_s(r). \quad (42)$$

Here, ε_{ab} is a constant polarization tensor satisfying (12) and (19). Now as before we focus on the off-diagonal mode and set $\varepsilon_{xy} = \varepsilon_{yx} = 1$ and $\varepsilon_{ab} = 0$ otherwise. So in this case, only non-zero perturbation will be given by

$$h_{xy} = \frac{r^2}{L^2} H_s(r) e^{-iM_s t}. \quad (43)$$

We further assume that $H_s(r) = H_{s0}(r) + H_{s1}(r)$ and insert (43) into the equation of motion for h_{xy} . Now we expand this differential equation in γ and find that zeroth order terms give us equation of motion for $H_{s0}(r)$, that is given by (21). The first order terms will give the equation of motion for $H_{s1}(r)$ and the full expression is given in the appendix A (equation (68)).

Similar to other screening masses, knowing the value of M_{s0} , we solve the coupled differential equations (68) and (21) for M_{s1} . The boundary condition for H_{s0} is that near the horizon, H_{s0} is given by (22) and it falls off as (9). $H_{s1}(r)$ also falls off as $1/r^4$ asymptotically. Also, we can use (22) and find its solution near the horizon, which is given by (69).

Given these boundary conditions and knowing the value of M_{s0} , we can solve the coupled differential equations iteratively for the correction M_{s1} . Finally we find that the screening mass for the spin-2 channel is given by

$$M_s = \pi T(3.404 - 167.619\gamma). \quad (44)$$

4 Comparison with QCD

In this section, we tabulate the results obtained using the procedures outlined in section 3 and appendix B for the screening masses and η/s . Using these results we would like to make a phenomenological study of certain properties of the QCD plasma. By adding higher derivative corrections we have introduced extra parameters in the theory. In keeping with the general strategy in [28], we will use input from lattice QCD for the energy density, the mass gap (and M_s) to fix these parameters. Using the fixed parameters, we will “predict” the value for η/s and M_v . Of course our results should be taken with a grain of salt since the corrections we have added are not the most general (since we have ignored covariant derivatives acting on the curvature tensor at six and eight derivative order). However, interestingly we will find that the predictions are reasonable giving hope that enlarging the space of couplings may eventually lead to sensible phenomenology. From the latest lattice calculations, the ratio of energy density with its free field limit for 4-dimensional lattice QCD at temperature $T \approx 2T_c$ is $(\varepsilon/\varepsilon_0)_{lattice} = 0.85 - 0.90$ [37, 38]. The values of screening masses for 2-flavour $N_f = 2$ QCD at $2T_c$ are : $(M)_{lattice}/\pi T = 1.68 - 1.91$; vector and spin-2 symmetry channels are $(M_v)_{lattice}/\pi T = 2.76 - 3.02$ and $(M_s)_{lattice}/\pi T = 2.48 - 2.64$ [16, 37, 39].

For comparison with lattice results, let us consider the following action with up to eight-derivative corrections

$$I = \frac{1}{2\ell_p^3} \int d^5x \sqrt{-g} \left[R + \frac{12}{L^2} + \alpha L^2 W_{abcd} W^{abcd} + \beta L^4 W^{abcd} W_{ac}{}^{ef} W_{bedf} + \gamma_i L^6 W_i^4 \right], \quad (45)$$

where α , β and γ_i are dimensionless coupling constants and they are systematically suppressed by the powers of ℓ_p^2/L^2 . In γ_i , i runs from 1 to 5 to give the general 8-derivative terms corresponding to table 2. This lagrangian will give a description of supersymmetric plasma if $\beta = 0$ and $\gamma_i W_i^4 = \gamma W_s^4$, where W_s^4 is given by (25). For a non-supersymmetric plasma, $\beta \neq 0$ and we also do not know the precise form of the $\gamma_i W_i^4$ term in this case.

Term	$M_v/\pi T$	$M_s/\pi T$
W^2	$4.322 + 12.965\alpha - 114.522\alpha^2$	$3.404 - 6.414\alpha + 128.620\alpha^2$
W^3	$4.322 + 45.738\beta$	$3.404 - 25.839\beta$
W_1^4	$4.322 + 603.164\gamma_1$	$3.404 + 279.749\gamma_1$
W_2^4	$4.322 + 238.119\gamma_2$	$3.404 - 532.634\gamma_2$
W_3^4	$4.322 - 285.403\gamma_3$	$3.404 - 122.103\gamma_3$
W_4^4	$4.322 - 695.936\gamma_4$	$3.404 - 307.493\gamma_4$
W_5^4	$4.322 - 216.641\gamma_5$	$3.404 - 673.803\gamma_5$

Table 3: Corrections to M_v, M_s

Implicitly in displaying these results we have assumed $O(\alpha^2) \sim O(\beta)$, $O(\alpha^3) \ll O(\beta)$, $O(\alpha^3) \ll O(\gamma_i)$ and $O(\beta^2) \ll O(\gamma_i)$. These are respected by the numerical solutions we have obtained in the discussion. Furthermore, we have included the α^2 terms to check the consistency of our numerics.

Let us start by assuming $\gamma_i W_i^4 = \gamma W_s^4$; we will soon comment on the general form. Now, from table 6, we can see that energy density for this lagrangian is given by

$$\varepsilon = \frac{3\pi^4 L^3 T^4}{2\ell_p^3} (1 + 18\alpha + 24\alpha^2 + 6\beta + 15\gamma). \quad (46)$$

Term	$4\pi\eta/s$	Mass gap/ πT
W^2	$(1 - 8\alpha + 112\alpha^2)$	$2.336 + 17.066\alpha - 97.704\alpha^2$
W^3	$(1 - 96\beta)$	$2.336 - 13.450\beta$
W_1^4	$(1 - 416\gamma_1)$	$2.336 + 266.317\gamma_1$
W_2^4	$(1 + 832\gamma_2)$	$2.336 - 532.634\gamma_2$
W_3^4	$(1 + 120\gamma_3)$	$2.336 - 139.515\gamma_3$
W_4^4	$(1 + 328\gamma_4)$	$2.336 - 272.673\gamma_4$
W_5^4	$(1 + 688\gamma_5)$	$2.336 - 691.218\gamma_5$

Table 4: Corrections to η/s and mass gap

We would like to take its ratio with the energy density of a non-supersymmetric conformal plasma in free field limit. To find energy density for free fields, we begin with the comparison of conformal anomalies of a four dimensional CFT with the one calculated using holographic techniques [28, 31]:

$$a = \pi^2 \frac{L^3}{\ell_p^3} \quad \text{and} \quad c = \pi^2 \frac{L^3}{\ell_p^3} (1 + 8\alpha). \quad (47)$$

Here a and c are central charges of CFT. Further, if we consider a free massless field theory with N_1 vectors, N_0 scalars and $N_{1/2}$ chiral fermions, we find that [29, 40]

$$a = \frac{124N_1 + 11N_{1/2} + 2N_0}{720}, \quad (48)$$

$$c = \frac{12N_1 + 3N_{1/2} + N_0}{120}, \quad (49)$$

$$t_4 = \frac{15(N_0 + 2N_1 - 2N_{1/2})}{2(N_0 + 12N_1 + 3N_{1/2})}. \quad (50)$$

Here t_4 is a constant that characterize the three-point function in CFT. This constant is found to be related to β by $t_4 = 4320\beta$. The original calculations for this term was done in [29] in absence of the quadratic terms and later it was argued in [31] that even in the presence of W_s^4 terms, the contribution to t_4 is of the form $\mathcal{O}(\beta\gamma, \gamma^2)$. So the expression for t_4 is correct so long we are considering only first order corrections from W^3 and W_s^4 terms. We are going to normalize by the energy density for a collection of free bosons and fermions which is given by

$$\varepsilon_0 = \frac{\pi^2 T^4}{30} \left(N_b + \frac{7}{8} N_f \right), \quad (51)$$

where N_b and N_f are the number of bosonic and fermionic degrees of freedom. For a non-supersymmetric theory $N_b \neq N_f$ where $N_b = N_0 + 2N_1$ and $N_f = 2N_{1/2}$. Now for the comparison of energy density (46) with its free field limit, we express quantities in (51) in terms of bulk gravity parameters using (47) - (51):

$$\varepsilon_0 = \frac{\pi^2 T^4}{30} \left(N_0 + 2N_1 + \frac{7}{4} N_{1/2} \right) = \frac{2\pi^4 L^3 T^4}{\ell_p^3} (1 + 16\alpha + 192\beta). \quad (52)$$

Now we can use (46), (52) and results from tables 3, 4, 6 to find that up to first order in β , γ and

second order in α

$$\frac{M}{\pi T} = 2.34 + 17.07\alpha - 97.70\alpha^2 - 13.45\beta - 139.51\gamma, \quad (53)$$

$$\frac{M_v}{\pi T} = 4.32 + 12.97\alpha - 114.52\alpha^2 + 45.74\beta - 394.35\gamma, \quad (54)$$

$$\frac{M_s}{\pi T} = 3.40 - 6.41\alpha + 128.62\alpha^2 - 25.84\beta - 167.62\gamma, \quad (55)$$

$$\frac{\varepsilon}{\varepsilon_0} = \frac{3}{4}(1 + 2\alpha - 8\alpha^2 - 186\beta + 15\gamma), \quad (56)$$

$$\frac{\eta}{s} = \frac{1}{4\pi}(1 - 8\alpha + 112\alpha^2 - 96\beta + 120\gamma). \quad (57)$$

To begin with, let us compare spectrum of type IIB theory with QCD to see what do we get. In this case, the supersymmetric plasma only has the W_s^4 correction and demanding $\epsilon/\epsilon_0 = [0.85, 0.90]$ leads⁵ to $\gamma = [0.008, 0.013]$. This leads to $\eta/s = [0.16, 0.20]$, $M/\pi T = [0.47, 1.10]$, $M_v/\pi T = [-0.94, 0.81]$ and $M_s/\pi T = [1.16, 1.91]$. Thus while η/s seems to be in the right ball park as RHIC and LHC data show [11], the screening masses are underestimated. In fact for the lower limit, M_v is also negative which indicates the possibility of not having a mass gap in the theory. Of course there is no reason for anything sensible with the leading W^4 correction which is relevant for the supersymmetric $\mathcal{N} = 4$ plasma. Thus we should enlarge the space of couplings. First, we could consider adding fundamental matter which corresponds to turning on the W^2 term. In this case, we can use the lattice free energy and mass gap results to fix γ, α to get $M_v/\pi T = [1.03, 2.40]$, $\eta/s \in [0.123, 0.154]$ and $M_s/\pi T \in [1.87, 2.46]$. This is already an improvement! However, note that certain values of α and γ , which produce $M_v < M$ or $M_s < M$, are not allowed as they violate our no level-crossing assumption. So let us now consider turning on W^3 and more general W^4 terms which would be relevant for a non-supersymmetric plasma to see what we get.

In Table 3, 4 and 6, we have mentioned the results for individual correction terms that might arise in the gravitational dual of non-supersymmetric field theories. Here we observe the following pattern: consider just the general set of W^4 terms. If the correction to the temperature is $T = \pi L^2/R_0(1 + c_i\gamma_i)$ then the corrected mass gap is $M/\pi T \approx (2.34 - 9.4c_i\gamma_i)$. Here γ_i is a dimensionless coupling constant and c_i depends on the different W_i^4 terms. Note that the above relation for the mass gap is approximate which we will use to get a range for the other physical quantities. Further, we also have the energy density $\varepsilon/\varepsilon_0 = 3/4(1 + c_i\gamma_i)$ and the screening mass for spin-2 channel is $M_s/\pi T \approx (3.40 - 9.3c_i\gamma_i)$. For the screening mass in the vector symmetry channel, we find that $M_v/\pi T = 4.32 + m_1c_i\gamma_i$ where m_1 varies between -4.2 and 24.0 for different cases. The corrections to η/s have similar qualitative behaviour and we find that the numerical coefficients of $c_i\gamma_i$ can vary between 8.0 and 14.9 . Now using these approximate results, we can reinstate the other corrections and compute the range for the coupling constants. To do that, similar to (53) - (57), we can write these physical quantities with contribution from all the W_i terms weighted by their respective coupling constants γ_i . In these expressions, we can define $\Gamma = c_i\gamma_i$. We replace $\gamma \rightarrow \Gamma/15$ in (53), (55) and (56). In (54) and (57), we replace $394.35\gamma \rightarrow m_1\Gamma$ and $120\gamma \rightarrow m_2\Gamma$. Here we have $m_1 \in [-4.2, 24.0]$ and $m_2 \in [8.0, 14.9]$. With these approximations, we have reduced the five coupling constants γ_i to only one, namely Γ . Now we can see that there are uncertainties in the screening mass for the vector channel because of m_1 , so we will compare M , M_s and $\varepsilon/\varepsilon_0$ with the lattice calculations. Then we calculate M_v and η/s , and compare these with lattice results.

⁵This value of γ would correspond to $\alpha_s = 0.14 - 0.20$ for $N_c = 3$.

By using the values of $(M)_{lattice}$, $(M_s)_{lattice}$ and $(\varepsilon/\varepsilon_0)_{lattice}$ in equations (53), (55) and (56) we get:

$$\alpha \in [0.00423, 0.02847], \quad \beta \in [-0.00060, 0.00003], \quad \Gamma \in [0.07300, 0.09179]. \quad (58)$$

Using (58) in (54) and (57), we find that⁶

$$\frac{\eta}{s} \in [0.11494, 0.19049], \quad (59)$$

$$\frac{M_v}{\pi T} \in [2.14224, 4.98317]. \quad (60)$$

Thus η/s is still well within⁷ experimental limits [11]. Interestingly, the range for screening mass in vector symmetry channel also encompasses the lattice results $(M_v)_{lattice}/\pi T = 2.76 - 3.02$. This gives hope that enlarging the space of couplings along the lines of [28, 31] may yield physically interesting results. In our numerics, we find that $O(\beta) \sim O(\alpha^2) \ll O(\gamma)$. Curiously, the dominant contribution comes from the γ -dependent terms.

5 Discussion

In this paper we have studied corrections to certain screening masses in the Yang-Mills plasma at strong coupling using AdS/CFT. On the gravity side the screening masses corresponded to the exchange of a graviton in the correlator of Polyakov loops. We were able to extract corrections to the mass gap which is the lightest supergravity mode. In addition we also computed corrections to the screening masses arising in the vector and spin-2 channels. We expanded the space of couplings by adding higher derivative corrections corresponding to the addition of fundamental matter and breaking of supersymmetry. Using lattice input for the mass gap and M_s , we were able to get sensible predictions for η/s and M_v .

Our approach may give the impression that all that one needs to do to find a realistic model is to add new couplings and adjust them to agree with observables available to us. It would be somewhat unsatisfying if that was all that could be said about this program. We would like to make the following interesting observation at this point. Recall that in order to get positive energy fluxes, we need to satisfy the following constraints for any CFT in 4 dimensions [29]:

$$1 - \frac{t_2}{3} - \frac{2}{15}t_4 \geq 0, \quad 1 + \frac{t_2}{6} - \frac{2}{15}t_4 \geq 0, \quad 1 + \frac{t_2}{3} + \frac{8}{15}t_4 \geq 0, \quad (61)$$

where $(c - a)/c = 8\alpha/(1 + 8\alpha) = t_2/6 + 4t_4/45$ for us. Remarkably, the ranges in (58) satisfy all these three inequalities. In other words, our findings are within the theoretical limits set forth in [29]. That this could happen was not at all guaranteed and we have examples of parameter spaces (e.g., considering the lower range for $M/\pi T$ to be 1.27) which would violate the above constraints.

Of course our numerical analysis should be taken with a grain of salt as we did not include the most general corrections possible. However, it will be a useful starting point for any future work on improving the AdS/CFT phenomenology program initiated in [28, 31]. For instance, it will be interesting to see the effect of adding a chemical potential to the corrections [41] or in the $\mathcal{N} = 2^*$ model [42].

⁶ Dropping the α^2 terms leads to a 10% change in the lower range of η/s and a 2% change in the upper range of M_v leaving everything else the same.

⁷Conservatively, for a RHIC/LHC plasma, $\eta/s < 0.4$.

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A Gauge conditions and other functions for eight derivative correction

For the eight derivative correction (25), the variables mentioned in (36) are:

$$c_1(r) = -H_1(r) \frac{2R_0^4}{r(3r^4 - R_0^4)} - H_0(r) \frac{2R_0^4}{L^6 r^{12} (3r^4 - R_0^4)^3} (675r^{20} + 648L^4 M_0^2 r^{14} R_0^4 - 225r^{16} R_0^4 - 1176L^4 M_0^2 r^{10} R_0^8 - 2808r^{12} R_0^8 + 632L^4 M_0^2 r^6 R_0^{12} + 876r^8 R_0^{12} - 104L^4 M_0^2 r^2 R_0^{16} + 1797r^4 R_0^{16} - 1275R_0^{20}) + \frac{dH_0(r)}{dr} \times \frac{8R_0^8(r^4 - R_0^4)(2L^4 M_0^2 r^2(3r^4 - R_0^4) - 3(36r^8 - 95r^4 R_0^4 + 49R_0^8))}{L^6 r^{11} (3r^4 - R_0^4)^2}, \quad (62)$$

$$d_1(r) = \frac{r^2}{2L^3 M_0^2} \frac{db_1(r)}{dr} - \frac{r^2 R_0^4}{(3r^4 - R_0^4) L^3 M_0^2} \frac{dH_1(r)}{dr} - \frac{12r^4 R_0^4}{(3r^4 - R_0^4)^2 L^2 M_0^2} H_1(r) + H_0(r) \frac{4R_0^4}{L^3 M_0^3 r^{11} (-3r^4 + R_0^4)^4} \left(-2L^8 M_0^5 r^4 R_0^4 (3r^4 - R_0^4)^3 + 6L^6 M_1 r^{16} (-3r^4 + R_0^4)^2 + L^4 M_0^3 r^2 R_0^4 (-1944r^{16} + 927r^{12} R_0^4 + 1041r^8 R_0^8 - 843r^4 R_0^{12} + 155R_0^{16}) - 9M_0(225r^{24} - 961r^{16} R_0^8 + 112r^{12} R_0^{12} + 839r^8 R_0^{16} - 80r^4 R_0^{20} - 135R_0^{24}) \right) - \frac{dH_0(r)}{dr} \frac{1}{2L^9 M_0^3 r^{10} (3r^4 - R_0^4)^3} \left(3L^6 M_1 r^{12} (3r^4 - R_0^4)^2 (r^4 + R_0^4) - 32L^4 M_0^3 r^2 R_0^8 (9r^{12} + 15r^8 R_0^4 - 21r^4 R_0^8 + 5R_0^{12}) + 18M_0 R_0^4 (75r^{20} + 551r^{16} R_0^4 - 1376r^{12} R_0^8 + 324r^8 R_0^{12} + 741r^4 R_0^{16} - 315R_0^{20}) \right). \quad (63)$$

The final equation of motion for $H_1(r)$ is following:

$$\begin{aligned}
& r(r^4 - R_0^4) \frac{d^2 H_1(r)}{dr^2} + (5r^4 - R_0^4) \frac{dH_1(r)}{dr} + \frac{r(64r^2 R_0^8 + L^4 M_0^2 (-3r^4 + R_0^4)^2)}{(3r^4 - R_0^4)^2} H_1(r) \\
&= \frac{H_0(r)}{L^6 r^{13} (3r^4 - R_0^4)^4} \left(-2L^{10} M_0 M_1 r^{14} (3r^4 - R_0^4)^4 - 48L^8 M_0^4 r^4 R_0^8 (3r^4 - R_0^4)^4 \right. \\
&\quad - 192R_0^8 (450r^{24} + 1803r^{20} R_0^4 - 9249r^{16} R_0^8 + 8288r^{12} R_0^{12} + 522r^8 R_0^{16} \\
&\quad - 2219r^4 R_0^{20} + 405R_0^{24}) - L^4 M_0^2 r^2 R_0^4 (6075r^{24} - 320841r^{20} R_0^4 + 952785r^{16} R_0^8 \\
&\quad - 1022070r^{12} R_0^{12} + 507657r^8 R_0^{16} - 117881r^4 R_0^{20} + 10291R_0^{24}) \Big) \\
&\quad - \frac{dH_0(r)}{dr} \frac{4R_0^4 (-r^4 + R_0^4)}{r^{12} (3r^4 - R_0^4)^3} \left(-2025r^{20} - 4680L^4 M_0^2 r^{14} R_0^4 + 51111r^{16} R_0^4 \right. \\
&\quad + 6024L^4 M_0^2 r^{10} R_0^8 - 169776r^{12} R_0^8 - 2328L^4 M_0^2 r^6 R_0^{12} + 209412r^8 R_0^{12} \\
&\quad \left. + 280L^4 M_0^2 r^2 R_0^{16} - 94623r^4 R_0^{16} + 12573R_0^{20} \right), \tag{64}
\end{aligned}$$

and the solution near the horizon is given by

$$\begin{aligned}
H_1(r) = & - \frac{M_0 L^3 (48L^4 M_0^3 + 1001M_0 R_0^2 + 2M_1 R_0^2) (L^5 M_0^2 + R_0^2 (16L + (16 + L^2 M_0^2) R_0))}{R_0^2 (L^4 M_0^2 + 16R_0^2)^2} \\
& + (r - R_0) \frac{M_0 L^3 (16 + M_0^2 L^2) (48L^4 M_0^3 + 1001M_0 R_0^2 + 2M_1 R_0^2)}{4R_0^2 (L^4 M_0^2 + 16R_0^2)}. \tag{65}
\end{aligned}$$

The equation of motion for $H_{v1}(r)$ is following:

$$\begin{aligned}
& r \frac{d^2 H_{v1}(r)}{dr^2} + 5 \frac{dH_{v1}(r)}{dr} + \frac{L^4 M_{v0}^2 r}{r^4 - R_0^4} H_{v1}(r) = - \frac{M_{v0} H_{v0}(r)}{L^2 r^{11} (r^4 - R_0^4)} \left(2L^6 M_{v1} r^{12} \right. \\
&\quad \left. + 48L^4 M_{v0}^3 r^2 R_0^8 + 15M_{v0} R_0^4 (5r^8 - 91r^4 R_0^4 + 85R_0^8) \right) \\
&\quad - \frac{dH_{v0}(r)}{dr} \frac{80R_0^8}{L^6 r^{12}} (8L^4 M_0^2 r^2 - 36r^4 + 27R_0^4), \tag{66}
\end{aligned}$$

and the solution near the horizon is given by

$$H_{v1}(r) = -(r - R_0) \frac{48L^8 M_0^4 + 2545L^4 M_0^2 R_0^2 + 2L^4 M_0 M_1 R_0^2 - 2880R_0^4}{L R_0^2 (L^4 M_0^2 + 20R_0^2)}. \tag{67}$$

The equation of motion for $H_{s1}(r)$ is given by

$$\begin{aligned}
& r(r^4 - R_0^4) \frac{d^2 H_{s1}(r)}{dr^2} + (5r^4 - R_0^4) \frac{dH_{s1}(r)}{dr} + rL^4 M_{s0}^2 H_{s1}(r) \\
&= - \frac{M_{s0} H_{s0}(r)}{L^2 r^{11}} \left(2L^6 M_{s1} r^{12} + 48L^4 M_{s0}^3 r^2 R_0^8 + M_{s0} R_0^4 (75r^8 - 789r^4 R_0^4 + 851R_0^8) \right) \\
&\quad - \frac{dH_{s0}(r)}{dr} \frac{12R_0^4 (r^4 - R_0^4)}{L^6 r^{12}} (25r^8 + 40L^4 M_{s0}^2 r^2 R_0^4 - 94r^4 R_0^4 - 129R_0^8), \tag{68}
\end{aligned}$$

and the solution near the horizon is

$$\begin{aligned}
H_{s1}(r) = & - \frac{2(48L^8 M_{s0}^5 + 2441L^4 M_{s0}^3 R_0^2 + 2L^{10} M_{s0}^2 M_{s1} R_0^2 + 4744M_{s0} R_0^4 + 16L^6 M_{s1} R_0^4)}{L^6 M_{s0} R_0^2 (L^4 M_{s0}^2 + 16R_0^2)} \\
& + (r - R_0) \frac{M_{s0}^2 (48L^8 M_{s0}^4 + 3977L^4 M_{s0}^2 R_0^2 + 2L^{10} M_{s0} M_{s1} R_0^2 + 7296R_0^4)}{4L^2 R_0^5 (L^4 M_{s0}^2 + 16R_0^2)}. \tag{69}
\end{aligned}$$

We display below the corrections to the geometry for the various higher derivative terms considered in this paper

Term	$f_1(r)$	$g_1(r)$
W^2	$-\frac{2R_0^4}{r^4} + \frac{8(50r^4R_0^4 - 17R_0^8)}{3r^8}\alpha$	$-\frac{2R_0^4}{r^4} + \frac{8(50r^4R_0^4 - 87R_0^8)}{3r^8}\alpha$
W^3	$\frac{-14r^4R_0^4 + 6R_0^8}{r^8}$	$\frac{2(-7r^4R_0^4 + 13R_0^8)}{r^8}$
W_1^4	$\frac{28R_0^4(5r^8 + 5r^4R_0^4 - 3R_0^8)}{r^{12}}$	$\frac{28R_0^4(5r^8 + 5r^4R_0^4 - 19R_0^8)}{r^{12}}$
W_2^4	$-\frac{56R_0^4(5r^8 + 5r^4R_0^4 - 3R_0^8)}{r^{12}}$	$-\frac{56R_0^4(5r^8 + 5r^4R_0^4 - 19R_0^8)}{r^{12}}$
W_3^4	$-\frac{15R_0^4(5r^8 + 5r^4R_0^4 - 3R_0^8)}{r^{12}}$	$-\frac{15R_0^4(5r^8 + 5r^4R_0^4 - 19R_0^8)}{r^{12}}$
W_4^4	$-\frac{29R_0^4(5r^8 + 5r^4R_0^4 - 3R_0^8)}{r^{12}}$	$-\frac{29R_0^4(5r^8 + 5r^4R_0^4 - 19R_0^8)}{r^{12}}$
W_5^4	$-\frac{74R_0^4(5r^8 + 5r^4R_0^4 - 3R_0^8)}{r^{12}}$	$-\frac{74R_0^4(5r^8 + 5r^4R_0^4 - 19R_0^8)}{r^{12}}$

Table 5: Corrections to metric

B Higher derivative corrections to η/s

In this appendix, we review the calculation of the shear viscosity of the holographic supersymmetric plasma with eight-derivative correction (25). We use the Kubo formula that relates the transport coefficients of the plasma to the field theory correlators and these correlators can be calculated using holographic techniques [43]. For higher derivative corrections we follow [13, 41].

Following the prescription set in [41], we use $u = R_0^2/r^2$ which is more convenient for the hydrodynamic calculations. The horizon is at $u = 1$ and the black hole solution for (24) is given by

$$ds^2 = -\frac{R_0^2}{L^2} \frac{f(u)}{u} dt^2 + \frac{L^2}{4u^2 g(u)} du^2 + \frac{R_0^2}{L^2} \frac{1}{u} (dx^2 + dy^2 + d\tau^2), \quad (70)$$

where

$$\begin{aligned} f(u) &= f_0(u)(1 + \gamma_1 f_1(u)), \\ g(u) &= f_0(u)(1 + \gamma_1 g_1(u)) \\ \text{and } f_0(r) &= (1 - u^2), \\ f_1(u) &= -15u^2(-5 - 5u^2 + 3u^4), \\ g_1(u) &= -15u^2(-5 - 5u^2 + 19u^4). \end{aligned} \quad (71)$$

Kubo formula relates the shear viscosity to the low frequency and zero momentum limit of the retarded Green's function of the stress tensor

$$\eta = -\lim_{\omega \rightarrow 0} \frac{1}{\omega} \text{Im} G_{xy,xy}^R(\omega, k=0), \quad (72)$$

where

$$G_{xy,xy}^R(\omega, k=0) = -i \int dt dx dy d\tau e^{i\omega t} \langle [T_{xy}(x), T_{xy}(0)] \rangle. \quad (73)$$

Now we translate the calculation of the correlator to the dual gravity by first calculating the effective action for the metric perturbation $h_x^y(t, u) = \int \frac{d^4 k}{(2\pi)^4} \phi_k(u) e^{-i\omega t + ik\tau}$. Evaluating the action (24) to quadratic order in the fluctuation $\phi_k(u)$ yields

$$I_\phi^{(2)} = \frac{1}{2\ell_p^3} \int \frac{d^4 k}{(2\pi)^4} du (A(u, \omega) \phi_k'' \phi_{-k} + B(u, \omega) \phi_k' \phi_{-k}' + C(u, \omega) \phi_k' \phi_{-k} + D(u, \omega) \phi_k \phi_{-k} + E(u) \phi_k'' \phi_{-k}'' + F(u) \phi_k'' \phi_{-k}' + \mathcal{K}), \quad (74)$$

where \mathcal{K} is the generalized Gibbons-Hawking boundary term and its detailed expression can be found in [41]. Now we follow the arguments given in [41] and directly read off the shear viscosity from the action (74)

$$\eta = \frac{1}{\ell_p^3} (\kappa_1(u) + \kappa_2(u))_{u=1}, \quad (75)$$

where

$$\begin{aligned} \kappa_1(u) &= \lim_{\omega \rightarrow 0} \sqrt{-\frac{g_{uu}(u)}{g_{tt}(u)}} \left(A(u, \omega) - B(u, \omega) + \frac{F'(u, \omega)}{2} \right), \\ \kappa_2(u) &= \lim_{\omega \rightarrow 0} \left(E(u, \omega) \left(\sqrt{-\frac{g_{uu}(u)}{g_{tt}(u)}} \right)' \right)'. \end{aligned} \quad (76)$$

Note that in (75) we have evaluated the quantities at the horizon whereas in (72) it was evaluated at the asymptotic boundary (at $u = 0$). For details of the arguments about these calculations, the reader is suggested to refer to [41]. The functions A, B, C, D, E, F at the horizon are given by

$$\begin{aligned} A(u, 0) &= \frac{2R_0^4(1-u^2)}{L^5 u} - \frac{10R_0^4 u(1-u^2)(15+15u^2-29u^4)\gamma}{L^5}, \\ B(u, 0) &= \frac{3R_0^4(1-u^2)}{2L^5 u} \\ &\quad + \frac{R_0^4 u(-45(5-16u^4+11u^6) + L^6(8u^2-52u^4+76u^6))\gamma}{2L^5}, \\ E(u, 0) &= \frac{16R_0^4 u^5(1-u^2)^2 \gamma}{L^3}, \\ F(u, 0) &= \frac{16R_0^4 u^4(1-3u^2+2u^4)\gamma}{L^4}. \end{aligned} \quad (77)$$

Using these we find that the shear viscosity is

$$\eta = \frac{\pi^3 L^3 T^3}{2\ell_p^3} (1 + 135\gamma). \quad (78)$$

The ratio of shear viscosity with entropy density (33) is

$$\frac{\eta}{s} = \frac{1}{4\pi} (1 + 120\gamma). \quad (79)$$

Term	$\left(\frac{\pi L^2}{R_0}\right) \times \text{Temperature}$	$\left(\frac{2l_p^3}{3\pi^4 L^3 T^4}\right) \times \text{Free energy}$	$\left(\frac{2l_p^3}{\pi^3 L^3 T^3}\right) \times \eta$
W^2	$\left(1 - 2\alpha - \frac{16}{3}\alpha^2\right)$	$(1 + 18\alpha + 24\alpha^2)$	$(1 + 10\alpha + 136\alpha^2)$
W^3	$(1 + 2\beta)$	$(1 + 6\beta)$	$(1 - 90\beta)$
W_1^4	$(1 - 28\gamma_1)$	$(1 - 28\gamma_1)$	$(1 - 444\gamma_1)$
W_2^4	$(1 + 56\gamma_2)$	$(1 + 56\gamma_2)$	$(1 + 888\gamma_2)$
W_3^4	$(1 + 15\gamma_3)$	$(1 + 15\gamma_3)$	$(1 + 135\gamma_3)$
W_4^4	$(1 + 29\gamma_4)$	$(1 + 29\gamma_4)$	$(1 + 357\gamma_4)$
W_5^4	$(1 + 74\gamma_5)$	$(1 + 74\gamma_5)$	$(1 + 762\gamma_5)$

Table 6: Corrections to thermodynamics and shear viscosity

The results for other correction terms are stated in Table 4 and 6. For the sake of completeness of the discussion, we have borrowed the results from [44] and [31] for the second order corrections to η .

References

- [1] J. M. Maldacena, “The large N limit of superconformal field theories and supergravity,” Adv. Theor. Math. Phys. **2**, 231 (1998) [Int. J. Theor. Phys. **38**, 1113 (1999)] [arXiv:hep-th/9711200].
- [2] O. Aharony, S. S. Gubser, J. M. Maldacena, H. Ooguri and Y. Oz, “Large N field theories, string theory and gravity,” Phys. Rept. **323**, 183 (2000) [arXiv:hep-th/9905111].
- [3] E. Witten, “Anti-de Sitter space and holography,” Adv. Theor. Math. Phys. **2**, 253 (1998) [arXiv:hep-th/9802150].
- [4] S. S. Gubser, I. R. Klebanov and A. M. Polyakov, “Gauge theory correlators from non-critical string theory,” Phys. Lett. B **428**, 105 (1998) [arXiv:hep-th/9802109].
- [5] E. V. Shuryak, “What RHIC experiments and theory tell us about properties of quark-gluon plasma?,” Nucl. Phys. A **750**, 64 (2005) [arXiv:hep-ph/0405066];
U. W. Heinz, “The strongly coupled quark-gluon plasma created at RHIC,” arXiv:0810.5529 [nucl-th].
- [6] C. V. Johnson and P. Steinberg, “What black holes teach about strongly coupled particles,” Phys. Today **63N5**, 29 (2010);
- [7] See, for example:
Y. V. Kovchegov, “AdS/CFT applications to relativistic heavy ion collisions: a brief review,” arXiv:1112.5403 [hep-ph].
S. Cremonini, “The Shear Viscosity to Entropy Ratio: A Status Report,” Mod. Phys. Lett. B **25**, 1867 (2011) [arXiv:1108.0677 [hep-th]].
J. Casalderrey-Solana, H. Liu, D. Mateos, K. Rajagopal and U. A. Wiedemann, “Gauge/String Duality, Hot QCD and Heavy Ion Collisions,” arXiv:1101.0618 [hep-th].
A. Sinha, “AdS/CFT and sQGP,” Nucl. Phys. A **862-863**, 168 (2011).
A. Sinha and R. C. Myers, “The Viscosity bound in string theory,” Nucl. Phys. A **830**, 295C

- (2009) [arXiv:0907.4798 [hep-th]].
- S. S. Gubser, “Using string theory to study the quark-gluon plasma: Progress and perils,” Nucl. Phys. A **830**, 657C (2009) [arXiv:0907.4808 [hep-th]].
- W. Zajc, “Quark Gluon Plasma at RHIC (and in QCD and String Theory),” presented at PASCOS 08 — see <http://pirsa.org/08060040/>;
- K. Rajagopal, “Quark Gluon Plasma in QCD, at RHIC, and in String Theory,” presented at PASCOS 08 — see <http://pirsa.org/08060041/>;
- D. T. Son and A. O. Starinets, “Viscosity, Black Holes, and Quantum Field Theory,” Ann. Rev. Nucl. Part. Sci. **57**, 95 (2007) [arXiv:0704.0240 [hep-th]].
- D. Mateos, “String Theory and Quantum Chromodynamics,” Class. Quant. Grav. **24** (2007) S713 [arXiv:0709.1523 [hep-th]];
- S. S. Gubser, “Heavy ion collisions and black hole dynamics,” Gen. Rel. Grav. **39** (2007) 1533 [Int. J. Mod. Phys. D **17** (2008) 673];
- D. T. Son, “Gauge-gravity duality and heavy-ion collisions,” AIP Conf. Proc. **957** (2007) 134.
- [8] P. Kovtun, D. T. Son and A. O. Starinets, “Viscosity in strongly interacting quantum field theories from black hole physics,” Phys. Rev. Lett. **94**, 111601 (2005) [arXiv:hep-th/0405231];
- [9] A. Buchel and J. T. Liu, “Universality of the shear viscosity in supergravity,” Phys. Rev. Lett. **93**, 090602 (2004) [hep-th/0311175].
- [10] T. Schafer and D. Teaney, “Nearly Perfect Fluidity: From Cold Atomic Gases to Hot Quark Gluon Plasmas,” Rept. Prog. Phys. **72**, 126001 (2009) [arXiv:0904.3107 [hep-ph]].
- C. Cao, E. Elliott, J. Joseph, H. Wu, J. Petricka, T. Schafer and J. E. Thomas, “Universal Quantum Viscosity in a Unitary Fermi Gas,” Science **331**, 58 (2011) [arXiv:1007.2625 [cond-mat.quant-gas]].
- M. Muller, J. Schmalian and L. Fritz, “Graphene: A nearly perfect fluid,” Phys. Rev. Lett. **103**, 025301 (2009).
- A. Sinha and B. Mukhopadhyay, “Implications of a viscosity bound on black hole accretion,” Phys. Lett. B **709**, 289 (2012) [arXiv:1108.5177 [astro-ph.HE]].
- [11] U. W. Heinz, C. Shen and H. Song, “The viscosity of quark-gluon plasma at RHIC and the LHC,” arXiv:1108.5323 [nucl-th].
- [12] J. Erdmenger, P. Kerner and H. Zeller, “Transport in Anisotropic Superfluids: A Holographic Description,” JHEP **1201**, 059 (2012) [arXiv:1110.0007 [hep-th]].
- A. Rebhan and D. Steineder, “Violation of the Holographic Viscosity Bound in a Strongly Coupled Anisotropic Plasma,” Phys. Rev. Lett. **108**, 021601 (2012) [arXiv:1110.6825 [hep-th]].
- [13] A. Buchel, J. T. Liu and A. O. Starinets, “Coupling constant dependence of the shear viscosity in N=4 supersymmetric Yang-Mills theory,” Nucl. Phys. B **707**, 56 (2005) [arXiv:hep-th/0406264].
- [14] A. Buchel, “Resolving disagreement for η/s in a CFT plasma at finite coupling,” Nucl. Phys. B **803**, 166 (2008) [arXiv:0805.2683 [hep-th]].
- [15] R. C. Myers, M. F. Paulos and A. Sinha, “Quantum corrections to η/s ,” Phys. Rev. D **79**, 041901 (2009) [arXiv:0806.2156 [hep-th]].

- [16] D. Bak, A. Karch and L. G. Yaffe, “Debye screening in strongly coupled $N=4$ supersymmetric Yang-Mills plasma,” JHEP **0708**, 049 (2007) [arXiv:0705.0994 [hep-th]].
- [17] C. Csaki, H. Ooguri, Y. Oz and J. Terning, “Glueball mass spectrum from supergravity,” JHEP **9901**, 017 (1999) [arXiv:hep-th/9806021].
- [18] J. A. Minahan, “Glueball mass spectra and other issues for supergravity duals of QCD models,” JHEP **9901**, 020 (1999) [arXiv:hep-th/9811156].
- [19] N. R. Constable and R. C. Myers, “Spin-two glueballs, positive energy theorems and the AdS/CFT correspondence,” JHEP **9910**, 037 (1999) [arXiv:hep-th/9908175].
- [20] R. C. Brower, S. D. Mathur and C. I. Tan, “Glueball Spectrum for QCD from AdS Supergravity Duality,” Nucl. Phys. B **587**, 249 (2000) [arXiv:hep-th/0003115].
- [21] M. J. Teper, “SU(N) gauge theories in 2+1 dimensions,” Phys. Rev. D **59**, 014512 (1999) [arXiv:hep-lat/9804008].
- [22] A. Hart, M. Laine and O. Philipsen, “Static correlation lengths in QCD at high temperatures and finite densities,” Nucl. Phys. B **586**, 443 (2000) [arXiv:hep-ph/0004060].
- [23] D. J. Gross and E. Witten, “Superstring Modifications Of Einstein’s Equations,” Nucl. Phys. B **277**, 1 (1986).
- [24] M. B. Green and C. Stahn, “D3-branes on the Coulomb branch and instantons,” JHEP **0309**, 052 (2003) [hep-th/0308061].
- [25] G. Policastro and D. Tsimpis, “ R^{*4} , purified,” Class. Quant. Grav. **23**, 4753 (2006) [hep-th/0603165].
- [26] A. Buchel, R. C. Myers, M. F. Paulos and A. Sinha, “Universal holographic hydrodynamics at finite coupling,” Phys. Lett. B **669**, 364 (2008) [arXiv:0808.1837 [hep-th]].
- [27] Y. Kats and P. Petrov, “Effect of curvature squared corrections in AdS on the viscosity of the dual gauge theory,” JHEP **0901**, 044 (2009) [arXiv:0712.0743 [hep-th]].
- [28] A. Buchel, R. C. Myers and A. Sinha, “Beyond $\eta/s = 1/4\pi$,” JHEP **0903**, 084 (2009) [arXiv:0812.2521 [hep-th]].
- [29] D. M. Hofman and J. Maldacena, “Conformal collider physics: Energy and charge correlations,” JHEP **0805**, 012 (2008) [arXiv:0803.1467 [hep-th]].
- [30] R. C. Myers and B. Robinson, “Black Holes in Quasi-topological Gravity,” JHEP **1008**, 067 (2010) [arXiv:1003.5357 [gr-qc]].
R. C. Myers, M. F. Paulos and A. Sinha, “Holographic studies of quasi-topological gravity,” JHEP **1008**, 035 (2010) [arXiv:1004.2055 [hep-th]].
- [31] A. Buchel, M. P. Heller and R. C. Myers, “sQGP as hCFT,” Phys. Lett. B **680**, 521 (2009) [arXiv:0908.2802 [hep-th]].
- [32] J. Noronha, M. Gyulassy and G. Torrieri, “Constraints on AdS/CFT Gravity Dual Models of Heavy Ion Collisions,” arXiv:0906.4099 [hep-ph].

- [33] P. K. Kovtun and A. O. Starinets, “Quasinormal modes and holography,” *Phys. Rev. D* **72**, 086009 (2005) [hep-th/0506184].
- [34] R. C. Myers, “Stress tensors and Casimir energies in the AdS/CFT correspondence,” *Phys. Rev. D* **60**, 046002 (1999) [arXiv:hep-th/9903203].
- [35] M. F. Paulos, “Higher derivative terms including the Ramond-Ramond five-form,” *JHEP* **0810**, 047 (2008) [arXiv:0804.0763 [hep-th]].
- [36] S. S. Gubser, I. R. Klebanov and A. A. Tseytlin, “Coupling constant dependence in the thermodynamics of $N = 4$ supersymmetric Yang-Mills theory,” *Nucl. Phys. B* **534**, 202 (1998) [arXiv:hep-th/9805156].
- [37] P. Petreczky, “Lattice QCD at non-zero temperature,” arXiv:1203.5320 [hep-lat].
- [38] M. Cheng *et al.*, “The QCD Equation of State with almost Physical Quark Masses,” *Phys. Rev. D* **77**, 014511 (2008) [arXiv:0710.0354 [hep-lat]];
A. Bazavov, T. Bhattacharya, M. Cheng, N. H. Christ, C. DeTar, S. Ejiri, S. Gottlieb and R. Gupta *et al.*, “Equation of state and QCD transition at finite temperature,” *Phys. Rev. D* **80**, 014504 (2009) [arXiv:0903.4379 [hep-lat]].
- [39] Y. Maezawa *et al.* [WHOT-QCD Collaboration], “Electric and Magnetic Screening Masses at Finite Temperature from Generalized Polyakov-Line Correlations in Two-flavor Lattice QCD,” *Phys. Rev. D* **81**, 091501 (2010) [arXiv:1003.1361 [hep-lat]].
S. Datta and S. Gupta, “Screening masses in SU(2) pure gauge theory,” *Phys. Lett. B* **471**, 382 (2000) [arXiv:hep-lat/9906023];
S. Datta and S. Gupta, “Dimensional reduction and screening masses in pure gauge theories at finite temperature,” *Nucl. Phys. B* **534**, 392 (1998) [arXiv:hep-lat/9806034];
A. Hart and O. Philipsen, “The spectrum of the three-dimensional adjoint Higgs model and hot SU(2) gauge theory,” *Nucl. Phys. B* **572**, 243 (2000) [arXiv:hep-lat/9908041].
- [40] See, for example N. D. Birrell and P. C. W. Davies, *Quantum Fields In Curved Space*, (Cambridge University Press, 1982).
- [41] R. C. Myers, M. F. Paulos and A. Sinha, “Holographic Hydrodynamics with a Chemical Potential,” *JHEP* **0906**, 006 (2009) [arXiv:0903.2834 [hep-th]].
S. Cremonini, K. Hanaki, J. T. Liu and P. Szepietowski, “Higher derivative effects on eta/s at finite chemical potential,” *Phys. Rev. D* **80**, 025002 (2009) [arXiv:0903.3244 [hep-th]].
- [42] C. Hoyos, S. Paik and L. G. Yaffe, “Screening in strongly coupled $N=2^*$ supersymmetric Yang-Mills plasma,” *JHEP* **1110**, 062 (2011) [arXiv:1108.2053 [hep-th]].
- [43] G. Policastro, D. T. Son and A. O. Starinets, “From AdS/CFT correspondence to hydrodynamics,” *JHEP* **0209**, 043 (2002) [arXiv:hep-th/0205052];
D. T. Son and A. O. Starinets, “Minkowski-space correlators in AdS/CFT correspondence: Recipe and applications,” *JHEP* **0209**, 042 (2002) [arXiv:hep-th/0205051].
- [44] N. Banerjee and S. Dutta, “Shear Viscosity to Entropy Density Ratio in Six Derivative Gravity,” *JHEP* **0907**, 024 (2009) [arXiv:0903.3925 [hep-th]].